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Neutrino Mass and the $SU(2)_R$ Breaking Scale

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Abstract

The neutrino sector in a left-right extension of the Standard Model depends on how $SU(2)_R$ is broken. I list all possible scenarios, including the ones where the Majorana ν_R mass is naturally much smaller than the $SU(2)_R$ breaking scale, which is desirable for generating the proper baryon asymmetry of the Universe. The best such choice is identified and discussed.

In the Standard Model of particle interactions, the neutrino is part of a left-handed doublet $(\nu, l)_L$ under $SU(2)_L \times U(1)_Y$. Whereas the charged lepton must have a right-handed singlet counterpart l_R, the singlet ν_R is not mandatory [because it is trivial under $SU(2)_L \times U(1)_Y$] and is absent in the minimal version of the model. On the other hand, its existence is usually assumed so that ν_L may acquire a naturally small Majorana mass as ν_R gets a large Majorana mass [again because it is trivial under $SU(2)_L \times U(1)_Y$] in the famous canonical seesaw mechanism [1, 2]. Where does ν_R come from? and what is the magnitude of its Majorana mass? The simplest answer [2] is that $U(1)_Y$ is actually a remnant of $SU(2)_R \times U(1)_{B-L}$ under which $(\nu, l)_R$ is a doublet, and the large Majorana ν_R mass comes from the vacuum expectation value (*vev*) of a scalar $SU(2)_R$ triplet, which also breaks $SU(2)_R \times U(1)_{B-L}$ to $U(1)_Y$. This scenario has dominated the thinking on neutrino mass for over 20 years, but it is not the only possibility, even if the existence of ν_R is conceded. (Mechanisms without ν_R are also possible and just as natural [3].) It may not even be the best possibility as far as leptogenesis [4] is concerned, because the $SU(2)_R$ gauge interactions will tend to diminish the ν_R number density in the early Universe.

Under $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$, the quarks and leptons transform as:

$$q_L = (u, d)_L \sim (3, 2, 1, 1/3), \tag{1}$$

$$q_R = (u, d)_R \sim (3, 1, 2, 1/3), \tag{2}$$

$$l_L = (\nu, e)_L \sim (1, 2, 1, -1), \tag{3}$$

$$l_R = (\nu, e)_R \sim (1, 1, 2, -1), \tag{4}$$

where the electric charge is given by

$$Q = I_{3L} + I_{3R} + \frac{1}{2}(B - L). \tag{5}$$

To break $SU(2)_R \times U(1)_{B-L}$ to $U(1)_Y$, there are two possibilities. One is to use the scalar

doublet

$$\Phi_R = (\phi_R^+, \phi_R^0) \sim (1,1,2,1), \tag{6}$$

the other is to use the scalar triplet

$$\xi_R = (\xi_R^{++}, \xi_R^+, \xi_R^0) \sim (1,1,3,2). \tag{7}$$

The subsequent breaking of $SU(2)_L \times U(1)_Y$ to $U(1)_{em}$ may be achieved with either a scalar doublet

$$\Phi_L = (\phi_L^+, \phi_L^0) \sim (1,2,1,1), \tag{8}$$

or a scalar bidoublet

$$\eta = \begin{pmatrix} \eta_1^0 & \eta_2^+ \\ \eta_1^- & \eta_2^0 \end{pmatrix} \sim (1,2,2,0), \tag{9}$$

where $I_{3L} = 1/2, -1/2$ for the rows, and $I_{3R} = -1/2, 1/2$ for the columns. The existence of a scalar triplet

$$\xi_L = (\xi_L^{++}, \xi_L^+, \xi_L^0) \sim (1,3,1,2) \tag{10}$$

may also be contemplated but its *vev* must be much smaller than that of Φ_L or η to be consistent with the precisely determined values of $\sin^2\theta_W$ and the masses of the W and Z bosons. Neutrino masses are sensitive to which of these 5 scalars are chosen, resulting in 5 basic scenarios, as described below.

(I) $\xi_R + \eta$

This is the canonical scenario where ν_L pairs up with ν_R through the *vev*'s of the bidoublet η to form a Dirac mass m_D while ν_R picks up a large Majorana mass m_R through the *vev* of the $SU(2)_R$ triplet ξ_R. The famous seesaw mass matrix

$$\mathcal{M}_\nu = \begin{pmatrix} 0 & m_D \\ m_D & m_R \end{pmatrix} \tag{11}$$

is obtained, with m_R of order the $SU(2)_R$ breaking scale. The zero of this matrix comes from the fact that there is no ξ_L.

(II) $\xi_R + \eta + \xi_L$

This is the canonical left-right symmetric scenario, where $\xi_L \leftrightarrow \xi_R$ is often imposed as a symmetry of the theory. Since the *vev* of ξ_L contributes to the Majorana ν_L mass, the neutrino mass matrix of Eq. (11) becomes

$$\mathcal{M}_\nu = \begin{pmatrix} m_L & m_D \\ m_D & m_R \end{pmatrix}. \tag{12}$$

This means that the canonical seesaw formula is corrected to read

$$m_\nu = m_L - \frac{m_D^2}{m_R}. \tag{13}$$

However m_L is routinely argued to be small because $\langle \xi_L^0 \rangle$ is of order $\langle \eta_1^0 \rangle \langle \eta_2^0 \rangle / \langle \xi_R^0 \rangle$ provided $m_{\xi_L}^2$ is positive and of order v_R^2. In this case, m_L may be larger or smaller than m_D^2/m_R, or the two terms may be of comparable magnitude.

For the many practitioners of the canonical seesaw mechanism, m_L is implicitly assumed to be negligible. On the other hand, if m_L is the dominant term, then ν_R may be dispensed with. In other words, we have just the Standard $SU(2)_L \times U(1)_Y$ Model with the simple addition of a Higgs triplet [5]. Again assuming $m_{\xi_L}^2$ to be positive and large, we have [6]

$$\langle \xi_L^0 \rangle = -\mu \langle \phi_L^0 \rangle^2 / m_{\xi_L}^2, \tag{14}$$

where μ is the $\xi_L^\dagger \Phi_L \Phi_L$ coupling. This mechanism without any ν_R is also a completely satisfactory explanation of the smallness of m_ν.

(III) $\Phi_R + \eta$

Here the *vev* of Φ_R breaks $SU(2)_R \times U(1)_{B-L}$ to $U(1)_Y$, and all fermions obtain Dirac masses from η. Since there is no ξ_R or ξ_L, the neutrino is apparently a Dirac particle in this scenario. Thus m_D has to be orders of magnitude smaller than any other Dirac mass. This is theoretically disfavored, and it is seldom discussed in the literature.

(IV) $\Phi_R + \eta + \Phi_L$

This is the left-right symmetric version of (III). Again the neutrino mass appears to be purely Dirac. However, the coexistence of Φ_L and η allows for an interesting extension of the usual left-right model, especially in the context of E_6. One of the complications of using a scalar bidoublet in a left-right extension of the Standard Model is that two different *vev*'s, i.e. $\langle \eta^0_{1,2} \rangle$, contribute to any given fermion mass, thus implying the existence of flavor changing neutral currents (FCNC) in the scalar sector [7]. This is not a problem if the $SU(2)_R$ breaking scale is very high as in models with a large Majorana ν_R mass. In models where the neutrino mass is purely Dirac, the $SU(2)_R$ breaking scale is not necessarily very high, so FCNC becomes the limiting constraint on the scale of $SU(2)_R$ breaking. This constraint may be relaxed if there exists [8] an exotic quark h of charge $-1/3$ such that $(u, h)_R$ is an $SU(2)_R$ doublet instead of the usual $(u, d)_R$. Then m_u comes from η^0_1, m_d comes from ϕ^0_L, and m_h comes from ϕ^0_R, with no FCNC in the scalar sector. This turns out to be a natural possibility [9] in the superstring-inspired E_6 model. As for the lepton sector, the Dirac mass partner of ν_L is then a new field which is a singlet, whereas the $SU(2)_R$ partner of e_R (usually called ν_R) is now a different particle. Because there are more neutral fermions in this extension, Majorana masses for ν_L may again be generated [10].

(V) $\Phi_R + \Phi_L$

This is the simplest way of breaking $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ to $U(1)_{em}$. However, since the bidoublet η is absent, there are apparently no fermion masses. On the other hand, this creates a unique opportunity, i.e. the possibility that all fermion masses, be they Dirac or Majorana, come from dimension-five operators instead [11, 12], i.e. operators suppressed by presumably the Planck mass. In the neutrino mass matrix of Eq. (12), m_L comes from $(l_L \Phi_L)^2$, m_R comes from $(l_R \Phi_R)^2$, and m_D comes from $(\bar{l}_L \Phi^\dagger_L)(l_R \Phi_R)$. The smallness of the Majorana neutrino mass compared to all Dirac fermion masses may then be attributed to

the smallness of $v_L \equiv \langle \phi_L^0 \rangle$ compared to $v_R \equiv \langle \phi_R^0 \rangle$.

There is another important consequence of this scenario. Because the Majorana ν_R mass is now given by v_R^2/Λ, where Λ may be of order the Planck mass, say 10^{19} GeV, it will be very much smaller than the $SU(2)_R$ breaking scale, i.e.

$$m_R \sim \frac{v_R^2}{\Lambda} << v_R. \tag{15}$$

This means that in the early Universe, at temperatures comparable to m_R, the $SU(2)_R$ gauge interactions of ν_R are strongly suppressed and can safely be ignored. This is a crucial requirement [13] for leptogenesis through the decay of ν_R [4]. Recent detailed analyses [14] of this mechanism for obtaining a realistic baryon asymmetry of the Universe and its relationship to the neutrino mass matrix all assume this implicitly.

Going back to Scenarios (III) and (IV), and allowing for the existence of $(l_L \Phi_L)^2$ in (IV) and that of $(l_R \Phi_R)^2$ in both (III) and (IV), the seesaw neutrino mass matrices of Eqs. (11) and (12) are again reproduced for (III) and (IV) respectively. This means that for a natural understanding of successful leptogenesis, the $SU(2)_R$ model to be adopted should be one with an $SU(2)_R$ doublet rather than a triplet. Scenario (V) requires v_R to be very high [11], of order the grand-unification scale, because of m_t. Scenario (IV) is a modification of (V) but without the v_R constraint, because fermion masses may now come from η. Scenario (III) is a special case of (IV) and has the desirable original form of the seesaw neutrino mass matrix, i.e. Eq. (11) and not Eq. (12) as in (IV) and (V). This previously neglected model should then be put forward as the model of choice for understanding both neutrino mass and leptogenesis.

The scalar sector of Scenario (III) consists of only Φ_R and η. Whereas $SU(2)_R \times U(1)_{B-L}$ is broken down to $U(1)_Y$ by the *vev* of the doublet Φ_R, $SU(2)_L \times U(1)_Y$ is broken down to $U(1)_{em}$ by the *vev*'s of the bidoublet η which also provide Dirac masses for all the fermions.

For example, consider the quark Yukawa couplings:

$$\mathcal{L}_Y = h_1 \bar{q}_L \eta q_R + h_2 \bar{q}_L \tilde{\eta} q_R + H.c., \tag{16}$$

where

$$\tilde{\eta} = \sigma_2 \eta^* \sigma_2 = \begin{pmatrix} \bar{\eta}_2^0 & -\eta_1^+ \\ -\eta_2^- & \bar{\eta}_1^0 \end{pmatrix}. \tag{17}$$

Hence

$$m_u = h_1 v_1 + h_2 v_2^*, \tag{18}$$

$$m_d = h_1 v_2 + h_2 v_1^*, \tag{19}$$

where $v_{1,2} \equiv \langle \eta_{1,2}^0 \rangle$. Because

$$Tr\tilde{\eta}^\dagger \tilde{\eta} = Tr\eta^\dagger \eta = \bar{\eta}_1^0 \eta_1^0 + \eta_1^+ \eta_1^- + \bar{\eta}_2^0 \eta_2^0 + \eta_2^+ \eta_2^-, \tag{20}$$

whereas the independent scalar quartic terms $f_1 \Phi_R^\dagger \tilde{\eta}^\dagger \tilde{\eta} \Phi_R$ and $f_2 \Phi_R^\dagger \eta^\dagger \eta \Phi_R$ contain $f_1 v_R^2 \bar{\eta}_1^0 \eta_1^0$ and $f_2 v_R^2 \bar{\eta}_2^0 \eta_2^0$ respectively, the effective scalar potential after Φ_R has been integrated out has different mass terms for η_1^0 and η_2^0, i.e.

$$(m^2 + f_1 v_R^2) \bar{\eta}_1^0 \eta_1^0 + (m^2 + f_2 v_R^2) \bar{\eta}_2^0 \eta_2^0. \tag{21}$$

This means that unless $f_1 = f_2$, it is impossible to make both coefficients negative and of order the electroweak breaking scale. In other words, either η_1^0 or η_2^0 must remain heavy, i.e. of order v_R. Let $m^2 + f_1 v_R^2 = -\mu^2$, then $m_2^2 = (f_2 - f_1) v_R^2 - \mu^2$, and v_2/v_1 is expected to be suppressed by a factor of order v_1^2/m_2^2, as shown for example in Ref.[15]. As a result, the contributions of v_2 to Eqs. (18) and (19) are negligible and the suppression of FCNC is achieved.

The model so far has only Dirac fermion masses. Majorana neutrino masses would normally come from the well-known dimension-five operator [16] $(l_L \Phi_L)^2$, but since Φ_L is absent, only $(l_R \Phi_R)^2$ is available. Thus the original seesaw matrix of Eq. (11) is obtained,

and m_R is guaranteed to be suppressed relative to v_R as shown in Eq. (15). Given the particle content of Scenario (III) and the acceptance of higher-dimensional operators, m_L actually gets a contribution from the dimension-seven operator $(\bar{l}_L \eta \tilde{\Phi}_R)^2/\Lambda^3$. However its magnitude is $v_1^2 v_R^2/\Lambda^3$ which is smaller than the double seesaw [11] mass of $v_1^2 \Lambda/v_R^2$ by the factor $(v_R/\Lambda)^4$.

Consider next the supersymmetric version of Scenario (III). Using the convention that all superfields are left-handed, q_R is replaced by $q^c \sim (3^*, 1, 2, -1/3)$ and l_R by $l^c \sim (1, 1, 2, 1)$. The Higgs sector now consists of the superfields η, Φ_R, and $\Phi_R^c \sim (1, 1, 2, -1)$. An extra unbroken discrete Z_2 symmetry is imposed, under which quark and lepton superfields are odd, but Higgs superfields are even. This serves to distinguish l^c from Φ_R, and leads to the usual R parity of most supersymmetric extensions of the Standard Model.

To break $SU(2)_R$ at a high scale without breaking the supersymmetry, consider the following superpotential:

$$W = M \epsilon_{ij} \Phi_{Ri} \Phi_{Rj}^c + \frac{1}{2\Lambda} (\epsilon_{ij} \Phi_{Ri} \Phi_{Rj}^c)^2. \tag{22}$$

Note that an extra nonrenormalizable term has been added. In this case, the scalar potential becomes

$$\begin{aligned} V &= \sum_j |M \epsilon_{ij} \Phi_{Ri} + \frac{1}{\Lambda} (\epsilon_{i'j'} \Phi_{Ri'} \Phi_{Rj'}^c) \epsilon_{ij} \Phi_{Ri}|^2 \\ &+ \sum_i |M \epsilon_{ij} \Phi_{Rj}^c + \frac{1}{\Lambda} (\epsilon_{i'j'} \Phi_{Ri'} \Phi_{Rj'}^c) \epsilon_{ij} \Phi_{Rj}^c|^2 + V_g, \end{aligned} \tag{23}$$

where V_g comes from the gauge interactions of Φ_R and Φ_R^c. Since the neutral components of Φ_R and Φ_R^c have opposite values of I_{3R} and $B - L$, the condition $V_g = 0$ at its minimum is satisfied if $\langle \Phi_R \rangle = \langle \Phi_R^c \rangle$ in the above. A supersymmetric vacuum ($V = 0$) is thus obtained with

$$v_R = \langle \Phi_R \rangle = \langle \Phi_R^c \rangle = \sqrt{M\Lambda}. \tag{24}$$

This shows that M of Eq. (22) may be identified with m_R of Eq. (15), i.e. the large Majorana mass of ν_R.

In this scenario, $SU(2)_R$ is broken at the scale v_R of Eq. (24). Below it, a consistent supersymmetric extension of the Standard Model survives, but with a singlet neutrino of Majorana mass $\sim M << v_R$. This singlet neutrino couples to $(\nu\eta_2^0 - e\eta_2^+)$ but its interaction with the $SU(2)_R$ gauge bosons are very much suppressed at the time of the early Universe when its temperature is comparable to M. Its decay will thus generate a lepton asymmetry [4, 14] which gets converted [17] into the present observed baryon asymmetry of the Universe through sphalerons during the electroweak phase transition.

So far the scale v_R has not been determined. There are two possible approaches. One is to assume that it has to do with gauge-coupling unification of the minimal supersymmetric standard model (MSSM) [18], in which case it should be 10^{16} GeV. This implies a singlet neutrino mass of order $(10^{16})^2/10^{19} = 10^{13}$ GeV. The other is to use present neutrino data [19, 20, 21] together with the requirement that the canonical seesaw matrix of Eq. (11) yields a satisfactory baryon asymmetry of the Universe through ν_R decay, from which a lower bound on v_R may be obtained. Recent indications [14] are that the smallest m_R is of order 10^8 GeV, which implies that v_R is at least of order 10^{14} GeV.

The bidoublet η contains two electroweak doublets and they are just right for the unification of gauge couplings in the MSSM. However, two such bidoublets are usually assumed in a supersymmetric model in order to have realistic quark and lepton masses. [Because of supersymmetry, we cannot use $\tilde{\eta}$ as the second bidoublet as in Eqs. (18) and (19).] In that case, there are four electroweak doublets and two would have to be heavy (i.e. at the v_R scale) not to spoil the unification of gauge couplings. An alternative possibility is to keep only one bidoublet and invoke a flavor-nondiagonal soft supersymmetry breaking scalar sector to account for the observed quark and lepton mass matrices [22].

Scenario (III) is distinguished by the absence of a Φ_L doublet. This is a problem if $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ is embedded in a larger symmetry such as $SO(10)$ or $[SU(3)]^3$, because any scalar multiplet of this larger symmetry would also contain Φ_L if it contains Φ_R. In that case, Eq. (12) is obtained, where m_L comes from $(l_L \Phi_L)^2$. As long as $m_L << m_D^2/m_R$, which holds if Λ is of order 10^{19} GeV, this is an acceptable scenario as well.

In conclusion, to understand both neutrino masses in terms of the original canonical seesaw mechanism, i.e. Eq. (11), and the success of leptogenesis through ν_R decay, the simplest and most natural model is $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ with a scalar sector consisting of only an $SU(2)_R$ doublet Φ_R and an $SU(2)_L \times SU(2)_R$ bidoublet η. The dimension-five operator $(l_R \Phi_R)^2/2\Lambda$ leads to a large Majorana mass for ν_R such that $m_R \simeq v_R^2/\Lambda << v_R$, which is desirable for generating the proper baryon asymmetry of the Universe. A successful supersymmetric version of this model has also been discussed.

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